

Chiral and $U(1)_A$ restorations high in the hadron spectrum, semiclassical approximation and large N_c .

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In quantum systems with large n (radial quantum number) or large angular momentum the semiclassical (WKB) approximation is valid. A physical content of the semiclassical approximation is that the quantum fluctuations effects are suppressed and vanish asymptotically. The chiral as well as $U(1)_A$ breakings in QCD result from quantum fluctuations. Hence these breakings must be suppressed high in the spectrum and the spectrum of high-lying hadrons must exhibit $U(2)_L \times U(2)_R$ symmetry of the classical QCD Lagrangian. This argument can be made stronger for mesons in the large N_c limit. In this limit all mesons are stable against strong decays and the spectrum is infinite. Hence, one can excite mesons of arbitrary large size with arbitrary large action, in which case the semiclassical limit is manifest. Actually we do not need the exact $N_c = \infty$ limit. For any large action there always exist such N_c that the isolated mesons with such an action do exist and can be described semiclassically. From the empirical fact that we observe multiplets of chiral and $U(1)_A$ groups high in the hadron spectrum it follows that $N_c = 3$ is large enough for this purpose.

1. Introduction

If one neglects the tiny masses of u and d quarks, which are much smaller than Λ_{QCD} , then the QCD Lagrangian exhibits the

$$U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A \quad (1)$$

symmetry. This is because the quark-gluon interaction Lagrangian in the chiral limit does not mix the left- and right-handed components of quarks and hence the total QCD Lagrangian for the two-flavor QCD can be split into the left-handed and right-handed parts which do not communicate with each other. We know that the $U(1)_A$ symmetry of the classical QCD Lagrangian is absent at the quantum level because of the $U(1)_A$ anomaly, which is an effect of quantum fluctuations [1]. We also know that the chiral $SU(2)_L \times SU(2)_R$ symmetry is spontaneously (dynamically) broken in the QCD vacuum [2]. That this is so is directly evidenced by the nonzero value of the quark condensate, $\langle \bar{q}q \rangle = \langle \bar{q}_L q_R + \bar{q}_R q_L \rangle \simeq -(240 \pm 10 MeV)^3$, which

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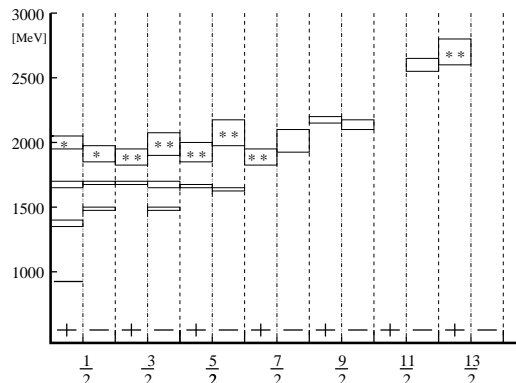


Figure 1. Excitation spectrum of the nucleon. The real part of the pole position is shown. Boxes represent experimental uncertainties. Those resonances which are not yet established are marked by two or one stars according to the PDG classification.

represents an order parameter for spontaneous chiral symmetry breaking. This quark condensate shows that in the QCD vacuum the left-handed quarks are correlated with the right-handed antiquarks (and vice versa) and hence the QCD vacuum breaks the chiral symmetry. This spontaneous (dynamical) breaking of chiral symmetry is a pure quantum effect based upon quantum fluctuations. To see the latter we remind the reader that the chiral symmetry breaking can be formulated via the Schwinger-Dyson (gap) equation. It is not yet clear at all which specific gluonic interactions are the most important ones as a kernel of the Schwinger-Dyson equation (e.g. instantons [3]², or gluonic exchanges [4], or perhaps other gluonic interactions, or a combination of different interactions). But in any case the quantum fluctuations of the quark and gluon fields are very strong in the low-lying hadrons and induce both chiral and $U(1)_A$ breakings. As a consequence we do not observe any chiral or $U(1)_A$ multiplets low in the hadron spectrum.

That the spontaneous breaking of chiral symmetry is an effect of quantum fluctuations of the quark field can be seen most generally from the definition of the quark condensate, which is a closed quark loop:

$$\langle \bar{\psi}\psi \rangle = -Tr \lim_{x \rightarrow 0_+} \langle 0 | T \{ \psi(0) \bar{\psi}(x) \} | 0 \rangle. \quad (2)$$

This closed quark loop explicitly contains a factor \hbar . The chiral symmetry breaking, which is necessarily a nonperturbative effect, is actually a (nonlocal) coupling of a quark line with the closed quark loop, which is a tadpole graph. Hence it always contains an extra factor \hbar as compared to the tree-level quark line.

The upper part of both baryon [5, 6] and meson [7, 8] spectra almost systematically exhibits multiplets of the chiral and $U(1)_A$ groups (for a pedagogical overview see [9]), though a careful experimental exploration of high-lying spectra must be done for a final conclusion. This phenomenon is referred to as effective chiral symmetry restoration or chiral symmetry restoration

²The instanton itself is an Euclidean semiclassical gluon field configuration. But chiral and $U(1)_A$ symmetry breakings by instantons is a quark field quantum fluctuations process. This is because the 't Hooft effective interaction is obtained only upon integrating of the $U(1)_A$ anomaly.

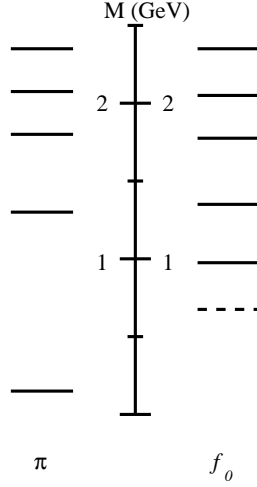


Figure 2. Pion and $n\bar{n}$ f_0 spectra.

of the second kind. This is illustrated in Fig. 1 and Fig. 2, where the excitation spectrum of the nucleon from the PDG compilation [16] as well as the excitation spectrum of π and f_0 (with the $\bar{n}n = \frac{\bar{u}u + \bar{d}d}{\sqrt{2}}$ content) mesons [7] are shown.³

Starting from the 1.7 GeV region the nucleon spectrum shows obvious signs of parity doubling. There are couple of examples where chiral partners of highly excited states have not yet been seen. Their experimental discovery would be an important task. All possible irreducible representations of the parity-chiral group necessarily contain parity doublets [6]. Similarly, in the chirally restored regime π and $\bar{n}n$ f_0 states must be systematically degenerate.

A quantitative measure of chiral symmetry breaking contribution to the hadron mass at the leading (linear) order is the *chiral asymmetry*

$$\chi = \frac{|M_1 - M_2|}{(M_1 + M_2)}, \quad (3)$$

where M_1 and M_2 are masses of particles within the same multiplet. This parameter has the interpretation of the part of the hadron mass due to chiral symmetry breaking. From the low-lying states the chiral asymmetry is typically 0.2 - 0.6, as can be seen e.g. from comparison of masses of the states $1/2^+$, $N(939)$ and $1/2^-$, $N(1535)$, π and σ , $\rho(770)$ and $a_1(1260)$, or $\rho(770)$ and $h_1(1170)$. If the chiral asymmetry is large as above, then it makes no sense to assign a given hadron to the chiral multiplet since its wave function is a strong mixture of different representations and we have to expect also large *nonlinear* symmetry breaking effects. Indeed, there is even no one-to-one mapping of the positive and negative parity hadrons low in the

³Those scalar mesons which are not $\bar{n}n$ states, and hence irrelevant for the present analysis, have been removed from the consideration. These are well established $f_0(1500)$ and $f_0(1710)$, which are believed to be mostly glueball and $\bar{s}s$ states, respectively, as well as $f_0(2102)$ which is seen in $\bar{p}p$ and considered by the authors of the partial wave analysis as a glueball due to its decay modes [21, 18].

Table 1

Chiral multiplets of π and $\bar{n}n - f_0$ mesons [7]. Comments: (i) $\pi(1300)$ and $f_0(1370)$ are well established states and can be found in the Meson Summary Table of Review of Particle Physics [16]. (ii) $\pi(1812 \pm 14)$ is abbreviated as $\pi(1800)$ in the Meson Summary Table of Review of Particle Physics [16]; $f_0(1770 \pm 12)$ is seen in $\bar{p}p$ as $\bar{n}n$ state [17, 18] and also recently as 15σ peak in the $\pi\pi$ channel in J/Ψ decays at ~ 1790 MeV [19]; The analysis of ref. [20] confirms a $\bar{n}n$ nature of this state. (iii) These states are clearly seen in a few different channels in $\bar{p}p$ [21, 18], though in order to appear in the Meson Summary Table of Review of Particle Physics they must be reconfirmed by independent experiment.

Chiral multiplet	Representation	χ	Spectral Overlap	Comment
$\pi(1300 \pm 100) - f_0(1370^{+130}_{-170})$	(1/2,1/2)	$0.03^{+0.09}_{-0.03}$	$0.1^{+0.3}_{-0.1}$	(i)
$\pi(1812 \pm 14) - f_0(1770 \pm 12)$	(1/2,1/2)	0.012 ± 0.007	0.09 ± 0.06	(ii)
$\pi(2070 \pm 35) - f_0(2040 \pm 38)$	(1/2,1/2)	$0.007^{+0.017}_{-0.007}$	$0.11^{+0.28}_{-0.11}$	(iii)
$\pi(2360 \pm 25) - f_0(2337 \pm 14)$	(1/2,1/2)	$0.005^{+0.009}_{-0.005}$	$0.08^{+0.13}_{-0.08}$	(iii)

spectrum. It indicates that here the chiral symmetry breaking effects are very strong and as a consequence the chiral symmetry is realized nonlinearly.

However, recent experimental data show that at meson masses about 1.8 - 2.3 GeV the chiral multiplets persist systematically (with a few missing states which should be discovered) and the chiral asymmetry is within 0.01 [7, 8].

A useful parameter that characterizes a "goodness" of symmetry in the spectrum is the *spectral overlap* which is defined as a ratio of the splitting within the multiplet to the distance between centers of gravity of two subsequent multiplets. Clearly, the symmetry is "good" and can be easily recognized if this parameter is much smaller than 1.

Both chiral asymmetries and spectral overlaps as well as assignments of mesons with $J = 0$ and $J = 2$, where the data sets are complete enough, are given in Tables 1 and 2.⁴

In the nucleon spectrum the chiral asymmetry is smaller than 0.02 for approximate parity doublets in the region $M \sim 1.7$ GeV. Similar doublets are observed in the delta spectrum starting from the $M \sim 1.9$ GeV. This means that the parity doubling in both cases is seen at approximately the same excitation energy with respect to the corresponding ground state. This indicates that not an absolute value of the energy is important in order to approach the chiral symmetry restoration regime but rather radial quantum number n or angular momentum J . Chiral multiplets in the nucleon spectrum are listed in Table 3.

In all these cases the hadrons can be believed to be members of chiral multiplets with a tiny admixture of other representations. It also means that practically the whole mass of the hadron *is not related* to the chiral symmetry breaking in the vacuum.

⁴Note that the Table 2 also evidences a restoration of $U(1)_A$ symmetry [7, 8, 9]. To see the latter one has to compare masses of states from the distinct (1/2,1/2) multiplets with the same isospin and opposite parity: $\pi_2 - a_2$, $f_2 - \eta_2$. The chiral asymmetries and spectral overlaps in these cases are similar to those ones in Table 2.

Table 2

Chiral multiplets of high-lying $J = 2$ mesons with the u,d valence quark content [8]. Comments: (i) These states are clearly seen in a few different channels in $\bar{p}p$ [21, 22, 18], though in order to appear in the Meson Summary Table of Review of Particle Physics they must be reconfirmed by independent experiment.

Chiral multiplet	Representation	χ	Spectral Overlap	Comment
$\omega_2(1975 \pm 20) - f_2(1934 \pm 20)$	(0,0)	0.01 ± 0.01	0.16 ± 0.15	(i)
$\omega_2(2195 \pm 30) - f_2(2240 \pm 15)$	(0,0)	0.01 ± 0.01	0.17 ± 0.17	(i)
$\pi_2(2005 \pm 15) - f_2(2001 \pm 10)$	(1/2,1/2)	$0.001^{+0.006}_{-0.001}$	$0.02^{+0.09}_{-0.02}$	(i)
$\pi_2(2245 \pm 60) - f_2(2293 \pm 13)$	(1/2,1/2)	$0.01^{+0.02}_{-0.01}$	$0.18^{+0.27}_{-0.18}$	(i)
$a_2(2030 \pm 20) - \eta_2(2030 \pm ?)$	(1/2,1/2)	$0.0 \pm ?$	$0.0 \pm ?$	(i)
$a_2(2255 \pm 20) - \eta_2(2267 \pm 14)$	(1/2,1/2)	$0.003^{+0.008}_{-0.003}$	$0.05^{+0.015}_{-0.05}$	(i)
$a_2(1950^{+30}_{-70}) - \rho_2(1940 \pm 40)$	$(0, 1) \oplus (1, 0)$	$0.003^{+0.018}_{-0.003}$	$0.04^{+0.27}_{-0.04}$	(i)
$a_2(2175 \pm 40) - \rho_2(2225 \pm 35)$	$(0, 1) \oplus (1, 0)$	$0.011^{+0.017}_{-0.011}$	$0.20^{+0.29}_{-0.20}$	(i)

Table 3

Chiral multiplets of excited nucleons [6]. Comments: (i) All these states are well established and can be found in the Baryon Summary Table of Review of Particle Physics. 200 MeV is taken as an interval between the consequent multiplets in order to evaluate the spectral overlap. (ii) There are two possibilities to assign the chiral representation: $(1/2, 0) \oplus (0, 1/2)$ or $(1/2, 1) \oplus (1, 1/2)$ because there is a possible chiral pair in the Δ spectrum with the same spin with similar mass.

Spin	Chiral multiplet	Representation	χ	Spectral Overlap	Comment
1/2	$N_+(1710 \pm 30) - N_-(1650^{+30}_{-10})$	$(1/2, 0) \oplus (0, 1/2)$	0.02 ± 0.02	0.3 ± 0.3	(i)
3/2	$N_+(1720^{+30}_{-70}) - N_-(1700^{+50}_{-50})$	$(1/2, 0) \oplus (0, 1/2)$	$0.01^{+0.03}_{-0.01}$	$0.1^{+0.4}_{-0.1}$	(i)
5/2	$N_+(1680^{+10}_{-5}) - N_-(1675^{+10}_{-5})$	$(1/2, 0) \oplus (0, 1/2)$	$0.002^{+0.006}_{-0.002}$	$0.025^{+0.1}_{-0.025}$	(i)
9/2	$N_+(2220^{+90}_{-40}) - N_-(2250^{+60}_{-80})$	see comment (ii)	$0.01^{+0.03}_{-0.01}$	$0.15^{+0.75}_{-0.15}$	(i),(ii)

2. Chiral symmetry restoration in excited hadrons by definition

By definition effective symmetry restoration means the following. In QCD hadrons with quantum numbers α are created when one applies the local interpolating field (current) $J_\alpha(x)$ with such quantum numbers on the vacuum $|0\rangle$. Then all hadrons that are created by the given interpolator appear as intermediate states in the two-point correlator

$$\Pi_{J_\alpha}(q) = i \int d^4x e^{-iqx} \langle 0 | T \{ J_\alpha(x) J_\alpha(0)^\dagger \} | 0 \rangle, \quad (4)$$

where all possible Lorentz and Dirac indices (specific for a given interpolating field) have been omitted. Consider two local interpolating fields $J_1(x)$ and $J_2(x)$ which are connected by a chiral transformation,

$$J_1(x) = U J_2(x) U^\dagger, \quad (5)$$

where

$$U \in SU(2)_L \times SU(2)_R \quad (6)$$

(or by the $U(1)_A$ transformation). Then if the vacuum was invariant under the chiral group,

$$U|0\rangle = |0\rangle, \quad (7)$$

it would follow from (4) that the spectra created by the operators $J_1(x)$ and $J_2(x)$ would be identical.

We know that in QCD one finds

$$U|0\rangle \neq |0\rangle. \quad (8)$$

As a consequence the spectra of the two operators must be in general different and we do not observe any chiral or $U(1)_A$ multiplets in the low-lying hadron spectra. However, it happens that the noninvariance of the vacuum becomes unimportant (irrelevant) high in the spectrum. Then the masses of the corresponding opposite parity hadrons, which are the members of the given parity-chiral multiplet, become close at large s (and identical asymptotically high),

$$M_1 - M_2 \rightarrow 0. \quad (9)$$

We stress that this effective chiral symmetry restoration does not mean that chiral symmetry breaking in the vacuum disappears, but that the role of the quark condensates that break chiral symmetry in the vacuum becomes progressively less important high in the spectrum. One could say that the valence quarks in high-lying hadrons *decouple* from the QCD vacuum.

3. Restoration of the classical symmetry in the semiclassical regime

In ref. [6] a justification for this chiral symmetry restoration has been suggested. Namely, at large space-like momenta $Q^2 = -q^2 > 0$ the correlator can be adequately represented by the operator product expansion, where all nonperturbative effects reside in different condensates [10]. The only effect that spontaneous breaking of chiral symmetry can have on the correlator is via the quark condensate of the vacuum, $\langle \bar{q}q \rangle$, and higher dimensional condensates that are not invariant under chiral transformation U . However, the contributions of all these condensates are suppressed by the inverse powers of momenta Q^2 . This shows that at large space-like momenta the correlation function becomes chirally symmetric. The dispersion relation provides a connection between the space-like and time-like domains of the correlator. In particular, the large Q^2 correlator is completely dominated by the large s spectral density $\rho(s)$, which is an observable. Hence the large s spectral density should be insensitive to the chiral symmetry breaking in the vacuum and the spectra of two operators $J_1(x)$ and $J_2(x)$ should approach each other,

$$\rho_1(s) - \rho_2(s) \rightarrow 0, \quad s \rightarrow \infty. \quad (10)$$

This is in contrast to the low s spectra which are very different because of the chiral symmetry breaking in the vacuum.

Unfortunately OPE at large Q^2 does not allow us to make quantitative statements concerning the functional behaviour that determines approaching the chiral-invariant regime at large s . This is because the OPE is only an asymptotic expansion and hence cannot be continued to the time-like region.

While the argument above on the asymptotic symmetry properties of spectral functions is rather robust (it is based actually only on the asymptotic freedom of QCD at large space-like momenta and on the analyticity of the two-point correlator), *a-priori* it is not clear whether it can be applied to the bound state systems, which the hadrons are. Indeed, it can happen that the asymptotic symmetry restoration applies only to that part of the spectrum, which is above the resonance region (i.e. where the current creates jets but not isolated hadrons). So the question arises whether it is possible to prove (or at least justify) the symmetry restoration in highly excited *isolated* hadrons. We show below that both chiral and $U(1)_A$ restorations in highly excited isolated hadrons can be anticipated as a direct consequence of the semiclassical regime in the highly excited hadrons, indeed.

At large n (radial quantum number) or at large angular momentum J we know that in quantum systems the *semiclassical* approximation (WKB) *must* work. Physically this approximation applies in these cases because the de Broglie wavelength of particles in the system is small in comparison with the scale that characterizes the given problem. In such a system as a hadron the scale is given by the hadron size while the wavelength of valence quarks is given by their momenta. Once we go high in the spectrum the size of hadrons increases as well as the typical momentum of valence quarks. This is why a highly excited hadron can be described semiclassically in terms of the underlying quark degrees of freedom.

The physical content of the semiclassical approximation is most transparently given by the path integral. The contribution of the given path to the path integral is regulated by the action $S(\phi(x))$ along the path $\phi(x)$ (the fields $\bar{\psi}, \psi, A$ are collectively denoted as ϕ) through the factor

$$\sim e^{i \frac{S(\phi(x))}{\hbar}}. \quad (11)$$

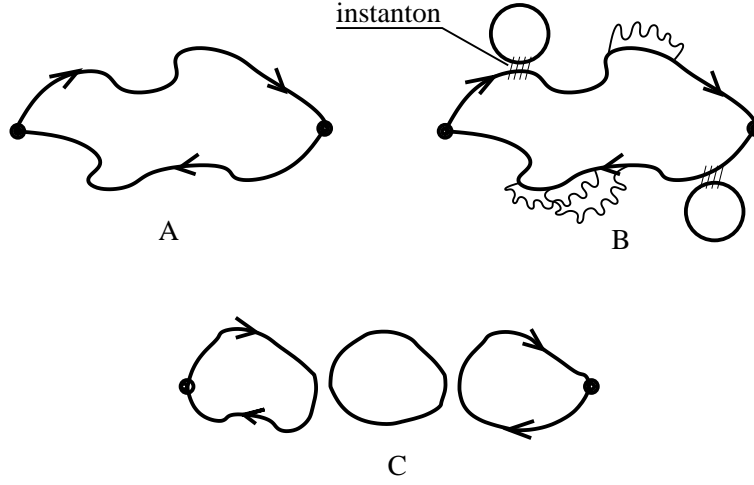


Figure 3. Tree level (A) and typical quantum fluctuations contributions (B and C) to the two-point function.

The semiclassical approximation applies when the action in the system $S \gg \hbar$. In this case the whole amplitude (path integral) is dominated by the classical path $\phi_{cl}(x)$ (stationary point) and those paths that are infinitesimally close to the classical path. All other paths that differ from the classical one by an appreciable amount do not contribute. These latter paths would represent the quantum fluctuation effects. In other words, in the semiclassical case the quantum fluctuations effects are strongly suppressed and vanish asymptotically. Mathematically it follows from the well known statement that when $\hbar \rightarrow 0$ the functional integral can be calculated in the stationary phase approximation (in Euclidean space it is a steepest descent or saddle point approximation). The stationary point is a solution of the classical equations of motion in the presence of the source $J(x)$,

$$\frac{\delta S}{\delta \phi(x)}[\phi_{cl}(J)] = -J(x). \quad (12)$$

Then the generating functional can be expanded in powers of \hbar as

$$W(J) = W_0(J) + \hbar W_1(J) + \dots, \quad (13)$$

where $W_0(J) = S(\phi_{cl}) + J\phi_{cl}$ and $W_1(J)$ represents contributions of the lowest order quantum fluctuations around the classical solution (determinant of the classical solution).⁵

The classical path, which is generated by W_0 , is a tree-level contribution to the path integral, see diagram a) on Fig. 3, and keeps chiral symmetries of the classical Lagrangian. Its contribution is of the order $(\hbar/S)^0$. The quantum fluctuations contribute at the orders $(\hbar/S)^1$ (the one loop order, generated by W_1), $(\hbar/S)^2$ (the two loops order), etc, see graphs b) and c) on Fig. 3.

⁵The expansion (13) should not be mixed with the quenched QCD. While there are no vacuum fermion loops in quenched QCD like those in Fig. 3c, there are still quantum fluctuations of the valence quark lines which induce chiral symmetry breaking, see Fig.3b. The first term in the expansion (13) does not contain any quantum fluctuations of the quark fields.

The $U(1)_A$ as well as the spontaneous $SU(2)_L \times SU(2)_R$ breakings result from quantum fluctuations of the quark fields and start from the one-loop order. However, in a hadron with large enough n or J , where action is large, the quantum fluctuations contributions must be relatively suppressed and vanish asymptotically. Then it follows that in such systems both the chiral and $U(1)_A$ symmetries must be restored. This is precisely what we see phenomenologically. In the nucleon spectrum the doubling appears either at large n excitations of baryons with a given small spin or in resonances of large spin. Similar features persist in the delta spectrum. In the meson spectrum the doubling is obvious for large n excitations of small spin mesons [7] and there are signs of doubling of large spin mesons (the data are, however, sparse). It would be certainly interesting and important to observe systematically multiplets of parity-chiral and parity- $U(1)_A$ groups (or, sometimes, when the chiral and $U(1)_A$ transformations connect *different* hadrons, the multiplets of the $U(2)_L \times U(2)_R$ group [7]). The high-lying hadron spectra must be systematically explored.

While the argument above is solid, theoretically it is not clear *a-priori* whether isolated hadrons still exist at excitation energies where a semiclassical regime is achieved. Hence it is conceptually important to demonstrate that in QCD hadrons still exist, while the dynamics inside such hadrons already is semiclassical. We do not know how to prove it for $N_c = 3$. However, the large N_c limit of QCD [11], while keeping all basic properties of QCD like asymptotic freedom and chiral symmetry, allows for a significant simplification. It is a good approximation, e.g., for the low-lying baryons [12].

In this limit it is known that all mesons represent narrow states, i.e. they are stable against strong decays. At the same time the spectrum of mesons is infinite (see, e.g., [13]). The latter is necessary to match the two-point function in the perturbation theory regime (which contains logarithm) at large space-like momenta with the discrete spectral sum in the dispersion integral. Then one can always excite a meson of any arbitrary large energy, which is of any arbitrary large size. In such a meson the action $S \gg \hbar$. Hence a description of this meson necessarily must be semiclassical. Then the equation of motion in such a meson must be according to some yet unknown solution of the classical QCD Lagrangian for a colorless meson. Hence it must exhibit chiral and $U(1)_A$ symmetries. This proves that it is possible to have an isolated hadron which can be described semiclassically.

Actually we do not need the exact $N_c = \infty$ limit for this statement. It can be formulated in the following way. For any large $S \gg \hbar$ there always exist such N_c that the isolated meson with such an action does exist and can be described semiclassically. From the empirical fact that we observe multiplets of chiral and $U(1)_A$ groups high in the hadron spectrum it follows that $N_c = 3$ is large enough for this purpose.

The strength of the argument given above is that it is very general. However we cannot say anything concrete about how all this happens. For that one needs a detailed microscopical understanding of both confinement and chiral symmetry breaking in QCD, and in particular a classical solution of the QCD Lagrangian for a highly-excited color-singlet hadron, around which an expansion (13) can be performed, which is a challenging task. But even though we can only assume how microscopically all this happens, it is a solid statement that in highly excited hadrons symmetries of the classical QCD Lagrangian should be observed. The only basis for this statement is that in such hadrons a semiclassical description is correct.⁶

⁶That the quantum fluctuations effects vanish in quantum bound state systems at large n or J is well known e.g. from the Lamb shift. The Lamb shift is a result of the radiative corrections (which represent effects of quantum

4. Conclusions

Using very general arguments we have shown that both $SU(2)_L \times SU(2)_R$ and $U(1)_A$ symmetries of the classical QCD Lagrangian should be approximately restored in highly excited hadrons and manifest asymptotically high. The reason is that in highly excited hadrons physics must necessarily be semiclassical. Unfortunately solutions of the classical QCD equations of motion for the white hadrons are yet unknown. Hence we can make only an assumption about the corresponding physical picture. If one constructs a highly excited hadron as a string (flux tube) with quarks at the ends that have definite chirality [14], which is a semiclassical picture, then one necessarily obtains all hadrons in chiral and $U(1)_A$ multiplets. Such a picture is rather natural and is well compatible with the Nambu string picture [15]. The ends of the string in the Nambu model move with velocity of light. Then, (it is an extension of the Nambu model) the quarks at the ends of the string must have definite chirality. In this way one is able to explain at the same time both Regge trajectories, chiral multiplet structure of excited hadrons and absence of the spin-orbit force in the u, d sector.

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fluctuations of electron and electromagnetic fields) and vanishes as $1/n^3$, and also very fast with increasing J . As a consequence high in the hydrogen spectrum the symmetry of the classical Coulomb potential gets restored. The author is grateful to D.O. Riska for suggesting this nice analogy. The other well-known example is the 't Hooft model [23] (QCD in 1+1 dimensions). In this model in the regime $N_c \rightarrow \infty$, $m_q \rightarrow 0$, $m_q \gg g \sim 1/\sqrt{N_c}$, the spectrum of the high-lying states is known exactly, $M_n^2 \sim n$. The chiral symmetry of the Lagrangian is broken (with no contradiction with the Coleman theorem since in this specific regime everything is determined by $N_c = \infty$, for any large but finite N_c the chiral symmetry is not broken in agreement with the Coleman theorem), which is due to gluon dressing of valence quarks and which is reflected by the fact that the positive and negative parity states are not degenerate and alternate in the spectrum. However, the mass difference between the neighbouring positive and negative parity states is $M_+ - M_- \sim 1/\sqrt{n}$ and vanishes high in the spectrum since the effect of quantum fluctuations dies out high in the spectrum. The latter can be explicitly seen from the fact that the amplitude of the higher quark Fock component in the meson wave function dies out very fast with increasing n [24]. The author is grateful to T. Cohen, T. DeGrand, S. Peris and A. Zhitnitsky for discussions on this point.

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